Constraints on the nonuniversal Z' couplings from $B \to \pi K$, πK^* and ρK Decays

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Abstract

Motivated by the large difference between the direct CP asymmetries $A_{CP}(B^- \to \pi^0 K^-)$ and $A_{CP}(\bar{B}^0 \to \pi^+ K^-)$, we combine the up-to-date experimental information on $B \to \pi K$, πK^* and ρK decays to pursue possible solutions with the nonuniversal Z' model. Detailed analyses of the relative impacts of different types of couplings are presented in four specific cases. Numerically, we find that the new coupling parameters, ξ^{LL} and ξ^{LR} with a common nontrivial new weak phase $\phi_L \sim -86^\circ$, which are relevant to the Z' contributions to the electroweak penguin sector $\triangle C_9$ and $\triangle C_7$, are crucial to the observed "πK puzzle". Furthermore, they are found to be definitely unequal and opposite in sign. We also find that $A_{CP}(B^- \to \rho^0 K^-)$ can put a strong constraint on the new Z' couplings, which implies the Z' contributions to the coefficient of QCD penguins operator O_3 involving the parameter ζ^{LL} required.

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1 Introduction

During the past several years, the observed discrepancies between the experimental measurements and the theoretical predications within the Standard Model (SM) for several observables in $B \to \pi K$ decays, the so-called "πK puzzle" [1], have attracted much attention. Extensive investigations both within the SM [2, 3, 4, 5, 6, 7], as well as with various specific New Physics (NP) scenarios [8, 9, 10], have been performed.

Averaging the recent experimental data from BABAR [11], Belle [12], CLEO [13] and CDF [14], the Heavy Flavor Averaging Group (HFAG) gives the following up-to-date rsults [15]

$$\begin{aligned} A_{CP}(B^- \to K^-\pi^0) &= 0.050 \pm 0.025 \ , \\ A_{CP}(\bar{B}^0 \to K^-\pi^+) &= -0.097 \pm 0.012, \end{aligned} \tag{1}$$

from which the difference between direct CP violations in the charged and the neutral modes

$$\Delta A \equiv A_{CP}(B^- \to K^-\pi^0) - A_{CP}(\bar{B}^0 \to K^-\pi^+) = 0.147 \pm 0.028 \tag{2}$$

is now established at about 5σ level.

Theoretically, it is generally expected that within the SM, these two CP asymmetries $A_{CP}(\bar{B}_d^0 \to \pi^+K^-)$ and $A_{CP}(B_u^- \to \pi^0K^-)$ should be approximately equal. For example, based on the QCD factorization approach (QCDF) [16], the recent theoretical predictions with two different schemes for the end-point divergence are

$$\begin{cases} A_{CP}(B_u^- \to \pi^0K^-) = -3.6\% \ , \\ A_{CP}(\bar{B}_d^0 \to \pi^+K^-) = -4.1\% \ , \end{cases} \quad \text{Scheme I (Scenario S4) [3]} \, , \tag{3}$$

$$\begin{cases} A_{CP}(B_u^- \to \pi^0K^-) = -10.8\% \ , \\ A_{CP}(\bar{B}_d^0 \to \pi^+K^-) = -12.4\% \ , \end{cases} \quad \text{Scheme II } (m_g = 0.5 \text{ MeV}) \text{ [8]} \, . \tag{4}$$

Here, the Scheme I is the way to parameterize the end-point divergence appearing in hard-spectator and annihilation corrections, by complex parameters $X_{A,H} = \int_0^1 dy/y = \ln(m_b/\Lambda)(1+\rho_{A,H}e^{i\phi_{A,H}})$, with $\rho_{A,H} \leq 1$ and unrestricted $\phi_{A,H}$ [3]. The Scheme II, as an alternative to the first one, is the way to quote the infrared finite gluon propagator to regulate the divergence. It is interesting to note that an infrared finite behavior of gluon propagator are not only obtained by solving the well-known Schwinger-Dyson equation [17, 18, 19], but also supported by recent Lattice QCD simulations [20]. However, both of these two schemes suffer the mismatch of ΔA given

by Eq. (2). Furthermore, within the framework of perturbative QCD approach (pQCD) [21], and the soft-collinear effective theory (SCET) [22], the theoretical predictions read

$$\begin{cases} A_{CP}(B_u^- \to \pi^0 K^-)_{PQCD} = (-1^{+3}_{-5})\% \;, \\ A_{CP}(\bar{B}_d^0 \to \pi^+ K^-)_{PQCD} = (-9^{+6}_{-8})\% \;, \end{cases} \quad \text{pQCD [5]}\,, \tag{5}$$

$$\begin{cases} A_{CP}(B_u^- \to \pi^0 K^-)_{SCET} = (-11 \pm 9 \pm 11 \pm 2)\% \;, \\ A_{CP}(\bar{B}_d^0 \to \pi^+ K^-)_{SCET} = (-6 \pm 5 \pm 6 \pm 2)\%. \end{cases} \quad \text{SCET [6]}\,. \tag{6}$$

Obviously, the present theoretical estimations within the SM are not consistent with the established ΔA. The mismatch might be due to our current limited understanding of the strong dynamics involved in hadronic B decays, but equally also to possible NP effects [23, 24].

In some well-motivated extensions of the SM, additional $U(1)'$ gauge symmetries and associated Z' gauge boson could arise. Searching for the extra Z' boson is an important mission in the experimental programs of Tevatron [25] and LHC [26]. Performing the constraints on the new Z' couplings through low-energy physics, on the other hand, is very imporatnt for the direct searches and understanding its phenomenology. Theoretically, the flavor changing neutral current (FCNC) is forbidden at tree level in the SM. One of the simple extensions is the family nonuniversal Z' model, which could be naturally derived in certain string constructions[27], E_6 models[28] and so on. It is interesting to note that the nonuniversal Z' couplings could lead to FCNC and new CP-violating effect [29], which possibly provide a solution to the afore mentioned "πK puzzle". With some simplifications of the nonuniversal Z' model and neglecting the color-suppressed electroweak (EW) penguins and the annihilation amplitudes, Ref. [9] gets four possible solutions

$$\begin{aligned} A_L: &\quad \{\xi^{LL}, \phi_L\} = \{0.0055, 110^\circ\}\,, \qquad B_L: \{\xi^{LL}, \phi_L\} = \{0.0098, -97^\circ\}\,, \quad \text{with}\, \xi^{LR} = 0\,; \\ A_{LR}: &\quad \{\xi^{LL} = \xi^{LR}, \phi_L\} = \{0.0104, -70^\circ\}\,, \qquad B_{LR}: \{\xi^{LL} = \xi^{LR}, \phi_L\} = \{0.0186, 83^\circ\}\,. \end{aligned} \tag{7}$$

However, the corresponding prediction $A_{CP}(B_u^- \to \pi^0 K^-) = -0.03 \pm 0.01$ [9] of solution A_L and A_{LR} in Eq. (7) is obviously inconsistent with the up-to-date experimental data 0.050 ± 0.025. Moreover, the annihilation amplitudes, which could generate some strong-interaction phases, are important for predicting CP violations.

Based on the above observations, in this paper we shall adopt the QCDF approach and reevaluate the effects of the nonuniversal Z' model on these decay modes with the updated

experimental data. Furthermore, since the $B \to \pi K^*$ and ρK decays also involve the same quark level $b \to s\bar{q}q$ ($q = u, d$) transition, it is necessary to take into account these decay modes.

In Section 2, we provide a quick survey of $B \to \pi K$, πK^* and ρK decays in the SM within the QCDF formalism; our numerical results, with two different schemes for the end-point divergence, are also presented. In Section 3, after reviewing the nonuniversal Z' model briefly, we present our analyses and numerical results in detail. Section 4 contains our conclusions. Appendix A recapitulates the decay amplitudes for the twelve decay modes within the SM [3]. Appendix B contains the formulas for hard-spectator and annihilation amplitudes with the infrared finite gluon propagator [8]. All the theoretical input parameters are summarized in Appendix C.

2 The SM results with two schemes for the end-point divergence.

In the SM, the effective Hamiltonian responsible for $b \to s$ transitions is given as [31]

$$\begin{aligned}\mathcal{H}_{\rm eff} \;=\; & \frac{G_F}{\sqrt{2}}\Big[V_{ub}V_{us}^*\left(C_1O_1^u + C_2O_2^u\right) + V_{cb}V_{cs}^*\left(C_1O_1^c + C_2O_2^c\right) - V_{tb}V_{ts}^*\big(\sum_{i=3}^{10}C_iO_i \\ & + C_{7\gamma}O_{7\gamma} + C_{8g}O_{8g}\big)\Big] + \text{h.c.},\end{aligned} \tag{8}$$

where $V_{qb}V_{qs}^*$ ($q = u$, c and t) are products of the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements [30], C_i the Wilson coefficients, and O_i the relevant four-quark operators whose explicit forms could be found, for example, in Refs. [2, 31].

In recent years, the QCDF approach has been employed extensively to study the hadronic B-meson decays. The $B \to \pi K$, πK^* and ρK decays have been studied comprehensively within the SM in Refs. [2, 3, 4, 32], and the relevant decay amplitudes within this formalism are shown in Appendix A. It is also noted that the framework contains estimates of the hard-spectator and annihilation corrections. Even though they are power-suppressed, their strength and associated strong-interaction phase are numerically important to evaluate the branching ratio and the CP asymmetry. However, unfortunately, the end-point singularities appear in twist-3 spectator and annihilation amplitudes. So, how to regulate the end-point divergence becomes important and necessary within this formalism. Here we shall adopt the following two schemes:

Scheme I: Parametrization

As the most popular way, the end-point divergent integrals are treated as signs of infrared sensitive contributions and phenomenologically parameterized by [2, 3]

$$\int_0^1 \frac{dy}{y} \to X_A = (1+\rho_A e^{i\phi_A}) \ln \frac{m_B}{\Lambda_h}, \qquad \int_0^1 dy \frac{\ln y}{y} \to -\frac{1}{2}(X_A)^2\,, \tag{9}$$

with $\Lambda_h = 0.5\,\mathrm{GeV}$, $\rho_A \leq 1$ and ϕ_A unrestricted. X_H is treated in the same manner. The different choices of ρ_A and ϕ_A correspond to different scenarios as discussed in Ref. [3], and S4 is mentioned as the most favorable one. It presents the moderate value of nonuniversal annihilation phase $\phi_A = -55^\circ$ (PP), -20° (PV) and -70° (VP). Conservatively, in our calculations we quote $\pm 5^\circ$ as their theoretical uncertainties. Taking $\rho_A = 1$ and $X_{A,H}$ universal for all decay processes belonging to the same modes (PP, PV or VP), we present our numerical results of branching ratios and direct CP asymmetries for $B \to \pi K$, πK^* and ρK decays in the third column of Tables 2 and 3, respectively.

As is known, the mixing-induced CP asymmetry A_{CP}^{mix} is well suited for testing the SM and searching for new physics effects. For example, the investigation of mixing-induced CP asymmetries in penguin dominated $\bar{B}^0 \to \pi^0 K_S^0$ and $\bar{B}^0 \to \rho^0 K_S^0$ decay modes has attracted much attention recently [33, 34, 35, 36]. After neglecting the $K_0 - \bar{K}_0$ mixing effect, the mixing-induced asymmetry could be written as

$$A_{CP}^{mix}(\bar{B}^0 \to f) = \frac{2\,\mathrm{Im}\lambda_f}{1+|\lambda_f|^2}\,, \qquad (f = \pi^0 K_S^0\,, \rho^0 K_S^0)\,, \tag{10}$$

with $\lambda_f = -\exp\{i\,\arg[\frac{V_{td}V_{tb}^*}{V_{td}^*V_{tb}}]\}\bar{A}_f/A_f$ in our phase convention. Our numerical predictions are listed in Table 4, which agree with the measurements within large experimental errors.

Scheme II: Infrared finite dynamical gluon propagator

In our previous paper [8], we have thoroughly studied the end-point divergence with an infrared finite dynamical gluon propagator. It is interesting to note that recent theoretical and phenomenological studies are now accumulating supports for a softer infrared behavior of the gluon propagator [19, 37, 38]. Furthermore, the infrared finite dynamical gluon propagator, which is shown to be not divergent as fast as $\frac{1}{q^2}$, has been successfully applied to the hadronic B-meson decays [39, 40]. In our evaluations, we shall quote the gluon propagator derived by

Cornwall (in Minkowski space) [17]

$$D(q^2) = \frac{1}{q^2 - M_g^2(q^2) + i\epsilon} \;, \tag{11}$$

where q is the gluon momentum. The corresponding strong coupling constant reads

$$\alpha_s(q^2) = \frac{4\pi}{\beta_0 \ln\left(\frac{q^2+4M_g^2(q^2)}{\Lambda_{QCD}^2}\right)} \;, \tag{12}$$

where $\beta_0 = 11 - \frac{2}{3}n_f$ is the first coefficient of the beta function, with n_f being the number of active quark flavors. The dynamical gluon mass square $M_g^2(q^2)$ is obtained as [17]

$$M_g^2(q^2) = m_g^2 \left[\frac{\ln\left(\frac{q^2+4m_g^2}{\Lambda_{QCD}^2}\right)}{\ln\left(\frac{4m_g^2}{\Lambda_{QCD}^2}\right)} \right]^{-\frac{12}{11}} \;, \tag{13}$$

where m_g is the effective gluon mass and $\Lambda_{QCD} = 225$ MeV. In Ref. [8], we present our suggestion, $m_g = 0.50 \pm 0.05$ GeV, which is a reasonable choice so that most of the observables (except for $A_{CP}(B \to \pi^0 K^-)$) are in good agreement with the experimental data. In this way, we find that the hard-spectator scattering contributions are real, and the annihilation contributions are complex with a large imaginary part [8]. Our numerical predictions for branching ratios, direct CP asymmetries and mixing-induced CP asymmetries are listed in the fourth column of Tables 2, 3 and 4, respectively.

Although numerically these two schemes have some differences, both of their predictions are consistent with most of the experimental data within errors. However, as expected in the SM, we again find that $A_{CP}(B_u^- \to \pi^0 K^-) = -0.041 \pm 0.008$ (-0.100 ± 0.008), are very close to $A_{CP}(\bar{B}_d^0 \to \pi^+ K^-) = -0.077 \pm 0.009$ (-0.116 ± 0.008) in the first (second) scheme. So, it is still hard to accommodate the measured large difference ΔA in the SM within the QCDF formalism, irrespective of adopting which scheme. In the following, we pursue possible solutions to this problem with a family nonuniversal Z' model [29].

3 Solution to the "πK puzzle" with nonuniversal Z' model.

3.1 Formalism of the family nonuniversal Z' model

A possible heavy Z' boson is predicted in many extensions of the SM, such as grand unified theories, superstring theories, and theories with large extra dimensions. The simplest way to

extend the SM gauge structure is to include a new $U(1)$ gauge group. A family nonuniversal Z' model can lead to FCNC processes even at tree level due to the non-diagonal chiral coupling matrix. The formalism of the model has been detailed in Ref. [29]. The relevant studies in the context of B physics have also been extensively performed in Refs. [9, 42, 43, 45].

After neglecting the $Z-Z'$ mixing with small mixing angle $\theta \sim \mathcal{O}(10^{-3})$ [44], and taking all the fields being the physical eigenstates, the Z' part of the neutral-current Lagrangian can be written as [29]

$$\mathcal{L}' = -g' J'_\mu Z'^\mu \,, \tag{14}$$

where g' is the gauge coupling constant of extra $U'(1)$ group at the EW M_W scale. The Z' chiral current is

$$J'_\mu = \bar{\psi}_i \gamma_\mu [(B_q^L)_{ij} P_L \,+\, (B_q^R)_{ij} P_R] \psi_j \,, \tag{15}$$

where ψ is the mass eigenstate of chiral fields and $P_{L,R} = (1 \mp \gamma_5)/2$. The effective chiral Z' coupling matrices are given as

$$B_q^X = V_{qX} \epsilon_{qX} V_{qX}^\dagger \,, \qquad (q = u, d; X = L, R) \,. \tag{16}$$

With the assumption of flavor-diagonal right-handed couplings , the Z' part of the effective Hamiltonian for $b \to s\bar{q}q \, (q = u, d)$ transitions can be written as [9]

$$\mathcal{H}_{eff}^{Z'} = \frac{2G_F}{\sqrt{2}} \big(\frac{g' M_Z}{g_1 M_{Z'}}\big)^2 \, B_{sb}^L (\bar{s}b)_{V-A} \sum_q \big(B_{qq}^L (\bar{q}q)_{V-A} + B_{qq}^R (\bar{q}q)_{V+A}\big) + h.c. \,, \tag{17}$$

where $g_1 = e/(\sin\theta_W \cos\theta_W)$ and $M_{Z'}$ the new gauge boson mass. It is noted that the forms of the above operators already exist in the SM. As a result, Eq. (17) can be modified as

$$\mathcal{H}_{eff}^{Z'} = -\frac{G_F}{\sqrt{2}} V_{ts}^* V_{tb} \sum_q (\Delta C_3 O_3^q + \Delta C_5 O_5^q + \Delta C_7 O_7^q + \Delta C_9 O_9^q) + h.c. \,, \tag{18}$$

where $O_i^q (i = 3, 5, 7, 9)$ are the effective operators in the SM, and ΔC_i the modifications to the corresponding SM Wilson coefficients caused by Z' boson, which are expressed as

$$\begin{aligned} \Delta C_{3,5} &= -\frac{2}{3 V_{ts}^* V_{tb}} \big(\frac{g' M_Z}{g_1 M_{Z'}}\big)^2 \, B_{sb}^L \, (B_{uu}^{L,R} + 2 B_{dd}^{L,R}) \,, \\ \Delta C_{9,7} &= -\frac{4}{3 V_{ts}^* V_{tb}} \big(\frac{g' M_Z}{g_1 M_{Z'}}\big)^2 \, B_{sb}^L \, (B_{uu}^{L,R} - B_{dd}^{L,R}) \,, \end{aligned} \tag{19}$$

in terms of the model parameters at the M_W scale.

Generally, the diagonal elements of the effective coupling matrices $B_{qq}^{L,R}$ are real as a result of the hermiticity of the effective Hamiltonian. However, the off-diagonal ones of B_{sb}^{L} can contain a new weak phase ϕ_L. Then, conveniently we can represent ΔC_i as[1]

$$\begin{aligned} \Delta C_{3,5} &= 2\,\frac{|V_{ts}^* V_{tb}|}{V_{ts}^* V_{tb}}\,\zeta^{LL,LR}\,e^{i\phi_L}\,, \\ \Delta C_{9,7} &= 4\,\frac{|V_{ts}^* V_{tb}|}{V_{ts}^* V_{tb}}\,\xi^{LL,LR}\,e^{i\phi_L}\,, \end{aligned} \tag{20}$$

where the real NP parameters $\zeta^{LL,LR}$, $\xi^{LL,LR}$ and ϕ_L are defined, respectively, as

$$\begin{aligned} \zeta^{LL,LR} &= -\frac{1}{3}\,\Big(\frac{g' M_Z}{g_1 M_{Z'}}\Big)^2\,\Big|\frac{B_{sb}^L}{V_{ts}^* V_{tb}}\Big|\,(B_{uu}^{L,R} + 2B_{dd}^{L,R})\,, \\ \xi^{LL,LR} &= -\frac{1}{3}\,\Big(\frac{g' M_Z}{g_1 M_{Z'}}\Big)^2\,\Big|\frac{B_{sb}^L}{V_{ts}^* V_{tb}}\Big|\,(B_{uu}^{L,R} - B_{dd}^{L,R})\,, \\ \phi_L &= \mathrm{Arg}[B_{sb}^L]\,. \end{aligned} \tag{21}$$

It is noted that the other SM Wilson coefficients may also receive contributions from the Z' boson through renormalization group (RG) evolution. With our assumption that no significant RG running effect between M_Z' and M_W scales, the RG evolution of the modified Wilson coefficients is exactly the same as the ones in the SM [31, 41]. For simplicity, we define

$$\begin{aligned} X' = \zeta^{LL} e^{i\phi_L}\,, \qquad Y' = \zeta^{LR} e^{i\phi_L}\,, \\ X = \xi^{LL} e^{i\phi_L}\,, \qquad Y = \xi^{LR} e^{i\phi_L}\,. \end{aligned} \tag{22}$$

The numerical results of Wilson coefficients in the naive dimensional regularization (NDR) scheme at the scale $\mu = m_b$ ($\mu_h = \sqrt{\Lambda_h m_b}$) are listed in Table 1. The values at the scale μ_h, with $m_b = 4.79$ GeV and $\Lambda_h = 500$ MeV, should be used in the calculation of hard-spectator and weak annihilation contributions.

3.2 Numerical analyses and discussions

With the theoretical formulas and the input parameters summarized in Appendix A, B and C, we now present our numerical analyses and discussions. Our analyses are divided into the following four cases with different simplifications for our attention, namely,

- Case I: With the simplifications $B_{uu}^{L,R} \simeq -2B_{dd}^{L,R}$ (*i.e.*, $\zeta^{LL,LR} = 0$) and $\xi^{LR} = 0$,

[1] For comparison, we take the same phase convention as Ref. [9].

Table 1: The Wilson coefficients C_i within the SM and with the contribution from Z' boson included in NDR scheme at the scale $\mu = m_b$ and $\mu_h = \sqrt{\Lambda_h m_b}$.

Wilson	$\mu = m_b$		$\mu_h = \sqrt{\Lambda_h m_b}$	
coefficients	C_i^{SM}	$\Delta C_i^{Z'}$	C_i^{SM}	$\Delta C_i^{Z'}$
C_1	1.075	$-0.006X$	1.166	$-0.008X$
C_2	-0.170	$-0.009X$	-0.336	$-0.014X$
C_3	0.013	$0.05X - 0.01Y - 2.20X' - 0.05Y'$	0.025	$0.11X - 0.02Y - 2.37X' - 0.12Y'$
C_4	-0.033	$-0.13X + 0.01Y + 0.55X' + 0.02Y'$	-0.057	$-0.24X + 0.02Y + 0.92X' + 0.09Y'$
C_5	0.008	$0.03X + 0.01Y - 0.06X' - 1.83Y'$	0.011	$0.03X + 0.02Y - 0.10X' + 0.09Y'$
C_6	-0.038	$-0.15X + 0.01Y + 0.1X' - 0.6Y'$	-0.076	$-0.32X + 0.04Y + 0.16X' - 1.26Y'$
C_7/α_{em}	-0.015	$4.18X - 473Y + 0.25X' + 1.27Y'$	-0.034	$5.7X - 459Y + 0.4X' + 1.7Y'$
C_8/α_{em}	0.045	$1.18X - 166Y + 0.01X' + 0.56Y'$	0.089	$3.2X - 355Y + 0.2X' + 1.5Y'$
C_9/α_{em}	-1.119	$-561X + 4.52Y - 0.8X' + 0.4Y'$	-1.228	$-611X + 6.7Y - 1.2X' + 0.6Y'$
C_{10}/α_{em}	0.190	$118X - 0.5Y + 0.2X' - 0.05Y'$	0.356	$207X - 1.4Y + 0.5X' - 0.1Y'$
$C_{7\gamma}$	-0.297	—	0.360	—
C_{8g}	-0.143	—	-0.168	—

- Case II: With the simplifications $B_{uu}^{L,R} \simeq -2B_{dd}^{L,R}$ only (*i.e.*, $\zeta^{LL,LR} = 0$),
- Case III: Taking $B_{uu}^{R} \simeq -2B_{dd}^{R}$ (*i.e.*, $\zeta^{LR} \simeq 0$), and leaving ζ^{LL} and $\xi^{LL,LR}$ arbitrary,
- Case IV: Without any simplifications for $B_{uu}^{L,R}$ and $B_{dd}^{L,R}$, *i.e.*, arbitrary values for $\zeta^{LL,LR}$ and $\xi^{LL,LR}$ are allowed.

Our fitting is performed with the experimental data varying randomly within their 2σ error-bars, while the theoretical uncertainties are obtained by varying the input parameters within the regions specified in Appendix C. In addition, we quote the Scheme II (taking $m_g = 0.5$GeV) to regulate the appearing end-point divergences.

With the assumption $B_{uu}^{L,R} \simeq -2B_{dd}^{L,R}$ and neglecting the color-suppressed EW penguins and the annihilation amplitudes, four possible solutions Eq. (7) to the "πK puzzle" are obtained in Ref. [9]. It is still worth to recheck these solutions with the updated experiment data and

Table 2: The CP-averaged branching ratios (in units of 10^{-6}) of $B \to \pi K$, πK^* and ρK decays in the SM with two end-point divergence regulation schemes, and in the nonuniversal Z' model with four different cases.

Decay Mode	Exp.	SM		Z' model			
	data	Scheme I	Scheme II	Case I	Case II	Case III	Case IV
$B_u^- \to \pi^- \overline{K}^0$	23.1 ± 1.0	19.0 ± 2.5	23.4 ± 3.9	23.3 ± 0.7	23.3 ± 0.6	23.2 ± 0.6	23.3 ± 0.7
$B_u^- \to \pi^0 K^-$	12.9 ± 0.6	10.5 ± 1.3	12.7 ± 2.0	12.5 ± 0.6	12.6 ± 0.6	12.5 ± 0.5	12.6 ± 0.6
$\overline{B}_d^0 \to \pi^+ K^-$	19.4 ± 0.6	16.2 ± 2.2	20.1 ± 3.4	19.9 ± 0.5	19.8 ± 0.5	19.9 ± 0.5	20.0 ± 0.5
$\overline{B}_d^0 \to \pi^0 \overline{K}^0$	9.8 ± 0.6	7.3 ± 1.1	9.3 ± 1.7	9.4 ± 0.6	9.5 ± 0.6	9.1 ± 0.4	9.1 ± 0.4
$B_u^- \to \pi^- \overline{K}^{*0}$	10.0 ± 0.8	11.7 ± 1.2	10.3 ± 3.3	8.4 ± 1.0	8.5 ± 0.9	8.7 ± 0.6	8.6 ± 0.7
$B_u^- \to \pi^0 K^{*-}$	6.9 ± 2.3	7.0 ± 0.7	6.0 ± 1.8	4.7 ± 0.6	4.7 ± 0.5	4.9 ± 0.3	4.8 ± 0.3
$\overline{B}_d^0 \to \pi^+ K^{*-}$	10.3 ± 1.1	9.9 ± 1.1	9.2 ± 2.8	7.5 ± 1.0	7.7 ± 0.9	8.0 ± 0.6	8.0 ± 0.6
$\overline{B}_d^0 \to \pi^0 \overline{K}^{*0}$	2.4 ± 0.7	4.1 ± 0.5	3.9 ± 1.3	3.6 ± 0.5	3.7 ± 0.4	3.5 ± 0.4	3.5 ± 0.4
$B_u^- \to \rho^- \overline{K}^0$	$8.0^{+1.5}_{-1.4}$	5.2 ± 0.9	10.6 ± 3.7	9.6 ± 1.4	9.7 ± 1.3	10.6 ± 1.3	10.7 ± 1.5
$B_u^- \to \rho^0 K^-$	$3.81^{+0.48}_{-0.46}$	2.5 ± 0.4	5.4 ± 1.6	4.22 ± 0.62	4.47 ± 0.63	4.7 ± 0.6	4.8 ± 0.7
$\overline{B}_d^0 \to \rho^+ K^-$	$8.6^{+0.9}_{-1.1}$	6.3 ± 1.0	13.0 ± 3.8	10.8 ± 1.4	10.9 ± 1.4	11.9 ± 1.4	12.5 ± 1.8
$\overline{B}_d^0 \to \rho^0 \overline{K}^0$	$5.4^{+0.9}_{-1.0}$	3.7 ± 0.5	7.3 ± 2.1	7.1 ± 0.9	7.4 ± 0.9	6.8 ± 0.9	6.9 ± 1.0

taken into account the neglected corrections. Furthermore, the possible solutions may also suffer strong constraints from $B \to \pi K^*$ and ρK decays, since they are also mediated by the same quark level $b \to s\bar{q}q$ transitions.

Case I: With the simplifications $B_{uu}^{L,R} \simeq -2B_{dd}^{L,R}(i.e.,\ \zeta^{LL,LR} = 0)$ and $\xi^{LR} = 0$

In this case, assuming $B_{uu}^{L,R} \simeq -2B_{dd}^{L,R}$ as in Ref. [9], the NP effect primarily manifests itself in the EW penguin sector and the Z' contribution to the Wilson coefficients Eq. (19) can be simplified as

$$\begin{aligned} \Delta C_{3,5} &= 0\,, \\ \Delta C_{9,7} &= 4\,\frac{|V_{ts}^* V_{tb}|}{V_{ts}^* V_{tb}}\,\xi^{LL,LR} e^{i\phi_L}\,, \quad \text{with}\, \xi^{LL,LR} = \Big(\frac{g' M_Z}{g_1 M_{Z'}}\Big)^2 \Big|\frac{B_{sb}^L}{V_{ts}^* V_{tb}}\Big|\, B_{dd}^{L,R}\,. \end{aligned} \tag{23}$$

As shown in Fig. 1 (a), taking $\xi^{LL} = 0.004$ and $\xi^{LR} = 0$, we find that $A_{CP}(B^- \to \pi^0 K^-)$ is

Table 3: The direct CP asymmetries (in unit of 10^{-2}) of $B \to \pi K$, πK^* and ρK decays. The other captions are the same as Table. 2.

Decay Mode	Exp.	SM		Z' model			
	data	Scheme I	Scheme II	Case I	Case II	Case III	Case IV
$B_u^- \to \pi^- \overline{K}^0$	0.9 ± 2.5	0.4 ± 0.1	0.04 ± 0.07	-1.6 ± 0.3	-2.7 ± 0.9	5.2 ± 0.5	5.1 ± 0.6
$B_u^- \to \pi^0 K^-$	5.0 ± 2.5	-4.1 ± 0.8	-10.0 ± 0.8	2.4 ± 1.6	2.3 ± 1.5	0.9 ± 0.7	1.2 ± 0.9
$\overline{B}_d^0 \to \pi^+ K^-$	$-9.8^{+1.2}_{-1.1}$	-7.7 ± 0.9	-11.6 ± 0.3	-11.7 ± 0.3	-11.0 ± 0.7	-10.5 ± 1.1	-10.5 ± 1.2
$\overline{B}_d^0 \to \pi^0 \overline{K}^0$	-1 ± 10	-1.5 ± 0.3	0.7 ± 0.3	-17 ± 2	-18 ± 2	-6 ± 2	-6 ± 2
$B_u^- \to \pi^- \overline{K}^{*0}$	$-2^{+6.7}_{-6.1}$	0.6 ± 0.1	0.09 ± 0.15	-2.1 ± 0.4	-3.3 ± 0.5	-0.6 ± 2.4	-3.0 ± 6.7
$B_u^- \to \pi^0 K^{*-}$	4 ± 29	-6 ± 2	-37 ± 9	6.8 ± 7.1	9.1 ± 7.2	-17 ± 4	-18 ± 6
$\overline{B}_d^0 \to \pi^+ K^{*-}$	-25 ± 11	-13 ± 2	-43 ± 10	-48 ± 3	-46 ± 3	-49 ± 3	-50 ± 5
$\overline{B}_d^0 \to \pi^0 \overline{K}^{*0}$	-15 ± 12	-4 ± 1	4 ± 2	-58 ± 9	-62 ± 9	-34 ± 7	-36 ± 11
$B_u^- \to \rho^- \overline{K}^0$	-12 ± 17	0.4 ± 0.1	0.5 ± 0.2	1.5 ± 0.1	-0.15 ± 0.7	5.3 ± 1.1	6.5 ± 4.5
$B_u^- \to \rho^0 K^-$	$41.9^{+8.1}_{-10.4}$	57.3 ± 5.8	42.3 ± 9.5	-36 ± 10	-46 ± 12	27 ± 4	27 ± 5
$\overline{B}_d^0 \to \rho^+ K^-$	15 ± 6	36 ± 4	29 ± 6	31 ± 3	33 ± 3	25 ± 2	25 ± 2
$\overline{B}_d^0 \to \rho^0 \overline{K}^0$	1 ± 20	-2.1 ± 1.3	-2.4 ± 1.4	45 ± 5	50 ± 5	8 ± 3	9 ± 4

Table 4: The mixing-induced CP asymmetries (in unit of 10^{-2}) of $\bar{B}^0 \to \pi^0 K_S^0$ and $\rho^0 K_S^0$ decays. The other captions are the same as Table. 2.

Decay Mode	Exp.	SM		Z' model			
	data	Scheme I	Scheme II	Case I	Case II	Case III	Case IV
$\overline{B}_d^0 \to \pi^0 K_S^0$	57 ± 17	77 ± 2	77 ± 2	46 ± 6	44 ± 6	61 ± 3	62 ± 5
$\overline{B}_d^0 \to \rho^0 K_S^0$	63^{+17}_{-21}	60 ± 2	66 ± 2	87 ± 2	84 ± 3	85 ± 3	86 ± 9

enhanced to be consistent with the experimental data when $\phi_L \sim -90^\circ$. Moreover, $A_{CP}(B^- \to \pi^- K^0)$ and $A_{CP}(B^0 \to \pi^+ K^-)$, which agree roughly with the experimental data in the SM, are not sensitive to the parameter ξ^{LL}. So, a possible solution to the observed "πK puzzle" Eq. (2) in Case I is naively favored.

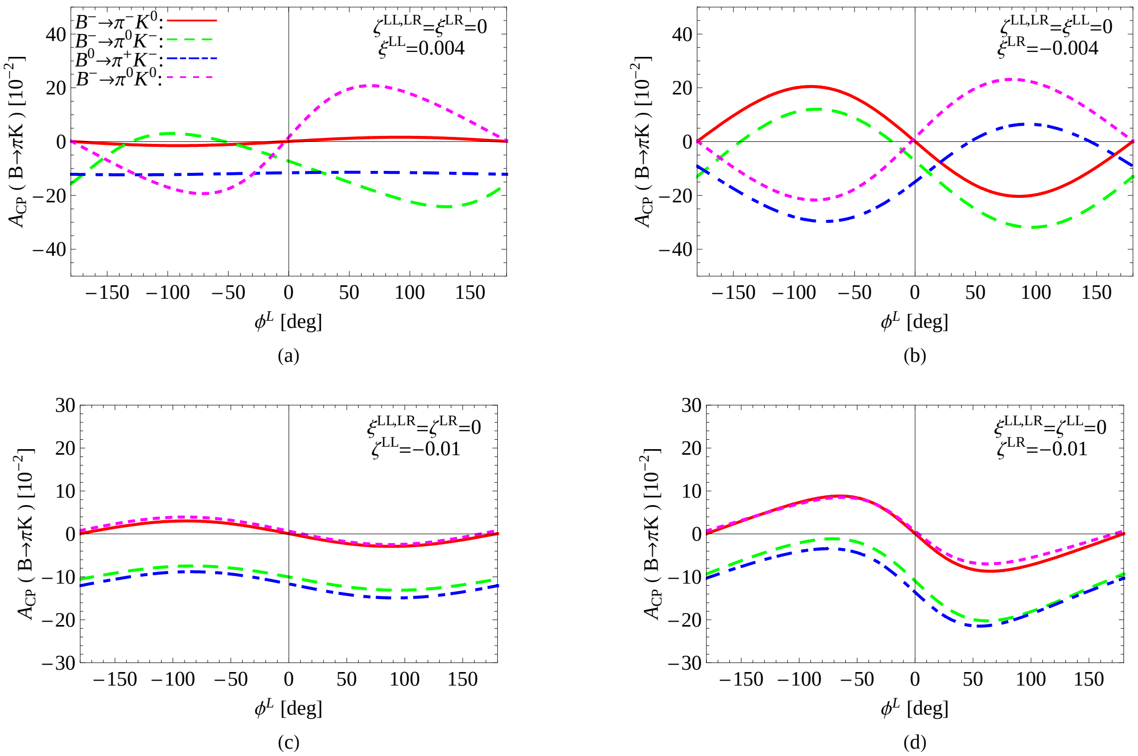

Figure 1: The dependence of $A_{CP}(B \to \pi K)$ on the new weak phase ϕ_L for the values of ξ^{LL}, ξ^{LR}, ζ^{LL}, and ζ^{LR} as marked by the legends.

Taking $\mathcal{B}(B \to \pi K)$ and $A_{CP}(B \to \pi K)$ as constraints on ξ^{LL} and ϕ_L, the allowed region for these two parameters are shown in Fig. 3 and the corresponding numerical results are listed in Table. 5, *i.e.*, $\xi^{LL} = (3.96 \pm 0.70) \times 10^{-3}$ and $\phi_L = -88^\circ \pm 7^\circ$. Our result confirms that the solution B_L in Eq. (7) is helpful to resolve the "πK puzzle" (note that a bit of difference might be due to the fact that the annihilation corrections are not included in Ref. [9]). However, the solution A_L is excluded by the updated experimental data $A_{CP}(B^- \to \pi^0 K^-) = 0.050 \pm 0.025$ as indicated in Fig. 1 (a).

With $\xi^{LL} = (3.96 \pm 0.70) \times 10^{-3}$ and $\phi_L = -88^\circ \pm 7^\circ$ as input parameters, we present our predictions for $\mathcal{B}(B \to \pi K^*, \rho K)$, $A_{CP}(B \to \pi K^*, \rho K)$ and $A_{CP}^{mix}(B^0 \to \pi^0 K_S, \rho^0 K_S)$ in the fifth column of Tables. 2, 3 and 4, respectively. We can see that most of them are consistent with the experimental data within 2σ. Especially, the predicted $A_{CP}^{mix}(B^0 \to \pi^0 K_S) = 0.46 \pm 0.06$ is very close to the measurement 0.57 ± 0.17 [15]. However, the prediction for $A_{CP}(B^- \to \rho^0 K^-) = -0.36 \pm 0.10$ presents a large discrepancy (larger than 6σ errors) with the current experiment data $0.419^{+0.081}_{-0.104}$ [15], which is also shown in Fig. 2 (a). This fact implies that

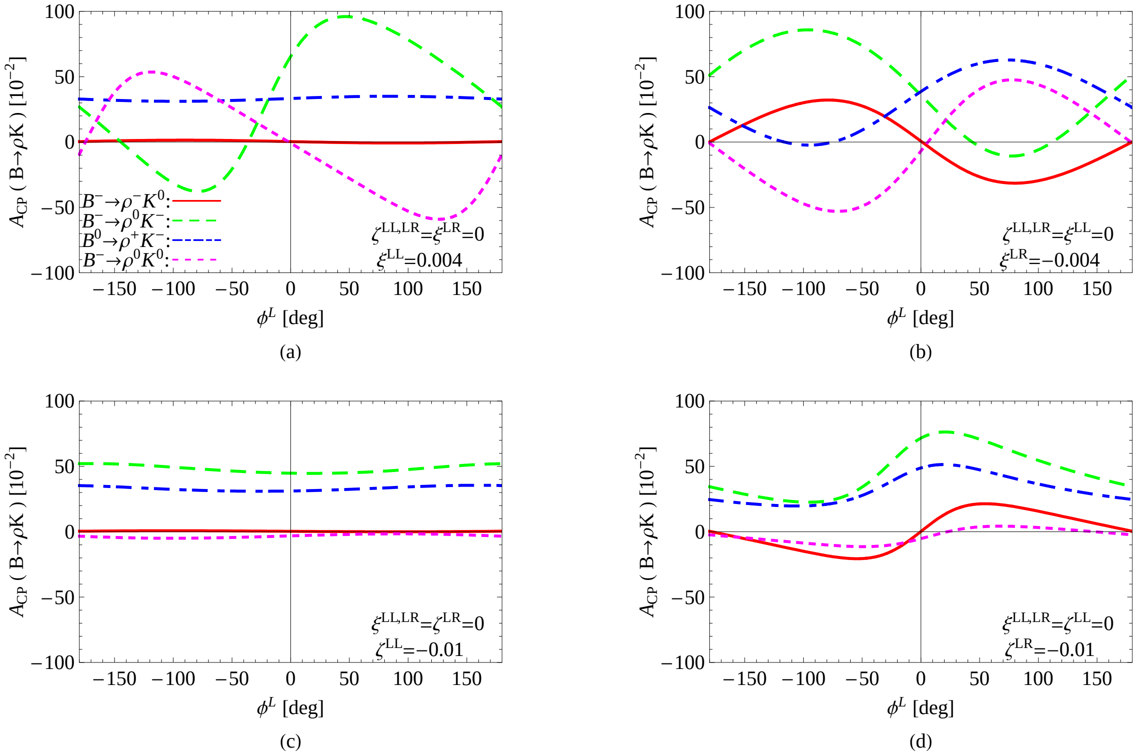

Figure 2: The dependence of $A_{CP}(B \to \rho K)$ on the new weak phase ϕ_L.

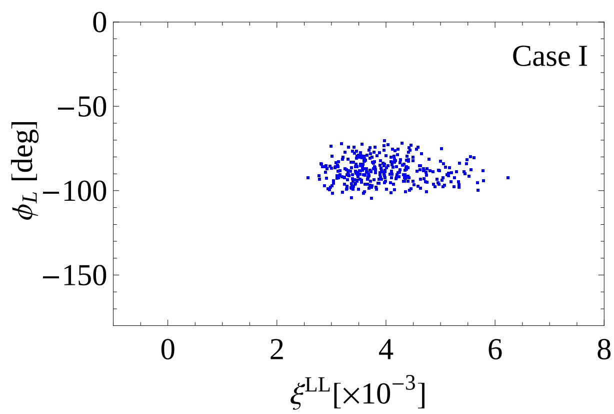

Figure 3: The allowed regions for the parameters ξ^{LL} and ϕ_L in Case I.

$A_{CP}(B^- \to \rho^0 K^-)$ can provide a strong constraint on the Z' couplings, at lease in Case I, and some more general Z' models might be required to explain all of these measurements.

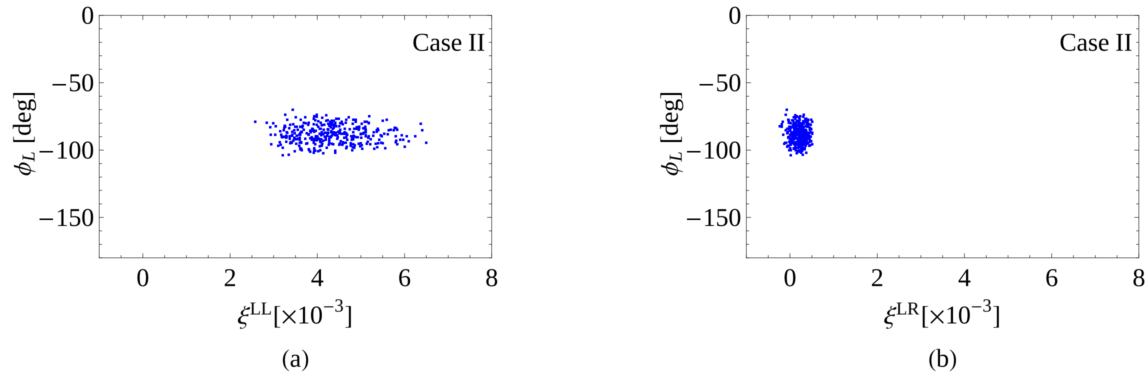

Figure 4: The allowed regions for the parameters $\xi^{LL,LR}$ and ϕ_L in Case II.

Case II: With the simplification $B_{uu}^{L,R} \simeq -2B_{dd}^{L,R}$ only ($i.\ e.\ \zeta^{LL,LR} = 0$).

It is interesting to note that, as shown in Fig. 1 (b), a region of minus ξ^{LR} with $\phi_L \sim -90^\circ$ can bridge the discrepancy of $A_{CP}(B^- \to \pi^0 K^-)$ between theoretical predictions and experimental data. Moreover, it is also possible to moderate the problem of $A_{CP}(B^- \to \rho^0 K^-)$ induced by ξ^{LL} as shown in Fig. 2 (b). So, in Case II we give up the simplification $\xi^{LR} = 0$ and pursue possible solutions to these discrepancies.

Taking $\mathcal{B}(B \to \pi K)$ and $A_{CP}(B \to \pi K)$ as constraints, we present the allowed regions for ξ^{LL}, ξ^{LR} and ϕ_L in Fig. 4. Unfortunately, we find that the required region of minus ξ^{LR} with $\phi_L \sim -90^\circ$ is excluded by $A_{CP}(B^0 \to \pi^+ K^-)$, because it will induce a large negative $A_{CP}(B^0 \to \pi^+ K^-)$ as shown in Fig. 1 (b). In addition, as shown in Fig. 1 (b), the region of plus ξ^{LR} with $\phi_L \sim -90^\circ$ is helpless to resolve the "πK puzzle". The Z' effects are therefore still dominated by large ξ^{LL}, and the problem of $A_{CP}(B^- \to \rho^0 K^-)$ induced by ξ^{LL} still exist.

In fact, with ξ^{LL} and ξ^{LR} having the same sign, the corresponding Z' contributions counteract with each other in the $B \to \pi^0 K^-$ decay as shown in Figs. 1 (a) and (b). It is also easily understood from the expression for the effective coefficient $\alpha_{3,EW}^p(PP) = a_9^p(PP) - a_7^p(PP)$ [3], which involves the leading-order Z' contribution in this case. Thus, we conclude that any attempt to explain the $B \to \pi K$ anomaly in the non-universal Z' model with the assumption $\xi^{LL} = \xi^{LR} = \xi$, as made in Ref. [45], is frangible and excluded in our case.

In a word, although the Z' contributions with a positive ξ^{LL} or a negative ξ^{LR} and $\phi_L \sim -90^\circ$ are helpful to bridge the discrepancy of $A_{CP}(B^- \to \pi^0 K^-)$, they would induce the unmatched

Table 5: The numerical results for the parameters $\xi^{LL,LR}$, $\zeta^{LL,LR}$ and ϕ_L in the four different cases. The dashes mean that the corresponding parameters are neglected in each case.

Parameters	Case I	Case II	Case III	Case IV
$\xi^{LL}(\times 10^{-3})$	3.96 ± 0.70	4.32 ± 0.75	1.52 ± 0.24	1.65 ± 0.35
$\xi^{LR}(\times 10^{-3})$	—	0.21 ± 0.15	-0.53 ± 0.13	-0.54 ± 0.15
$\zeta^{LL}(\times 10^{-3})$	—	—	-11.8 ± 3.1	-14.6 ± 7.1
$\zeta^{LR}(\times 10^{-3})$	—	—	—	1.04 ± 2.70
ϕ^L	$-88^\circ \pm 7^\circ$	$-88^\circ \pm 7^\circ$	$-86^\circ \pm 14^\circ$	$-85^\circ \pm 16^\circ$

$A_{CP}(B^- \to \rho^0 K^-)$ and $A_{CP}(B^0 \to \pi^+ K^-)$, respectively. Thus, with both $\mathcal{B}(B \to \pi K)$ and $A_{CP}(B \to \pi K, \rho K)$ as constraints, our results indicate that all of the parameter spaces in Case I and Case II are excluded with the assumption $B^{L,R}_{uu} \simeq -2B^{L,R}_{dd}$. As an alternative, in the following, we proceed to pursue possible solutions to these observations by considering the Z' contributions to the QCD penguins $\triangle C_{3,5}$.

Case III: Taking $B^R_{uu} \simeq -2B^R_{dd}$($i.e.$, $\zeta^{LR} \simeq 0$), and leaving ζ^{LL} and $\xi^{LL,LR}$ arbitrary.

As shown in Fig. 1 (c), we find that the variation trends of $A_{CP}(B^0 \to \pi^+ K^-)$ and $A_{CP}(B^- \to \pi^0 K^-)$ are always the same, indicating that the Z' contributions in this case could not give a solution to the observed "πK puzzle" directly, as well as the unmatched $A_{CP}(B^- \to \rho^0 K^-)$ induced by ξ^{LL}. However, it is interesting to note that, with $\phi_L \sim -90^\circ$, both $A_{CP}(B^0 \to \pi^+ K^-)$ and $A_{CP}(B^- \to \pi^0 K^-)$ could be enhanced simultaneously, which may relax the constraints on ξ^{LR}. As mentioned in Case II, a negative ξ^{LR} is favored by the "πK puzzle" and can moderate the problem of $A_{CP}(B^- \to \rho^0 K^-)$ induced by ξ^{LL}. So, the parameter ζ^{LL} may play an important role.

With $\mathcal{B}(B \to \pi K)$, $A_{CP}(B \to \pi K)$ and $A_{CP}(B \to \rho K)$ as constraints, the allowed regions for ξ^{LL}, ξ^{LR}, ζ^{LL} and ϕ_L are shown in Figs. 5. We find that none of ξ^{LL}, ξ^{LR} and ζ^{LL} could be neglected. Especially, the ζ^{LL} part moderates the contradictions caused by ξ^{LL} and ξ^{LR}. Furthermore, it is interesting to note that our predictions for $\mathcal{B}(B \to \pi K^*, \rho K)$, $A_{CP}(B \to \pi K^*)$ and $A^{mix}_{CP}(B^0 \to \pi^0 K_S, \rho^0 K_S)$, listed in Tables 2, 3 and 4, respectively, are all consistent

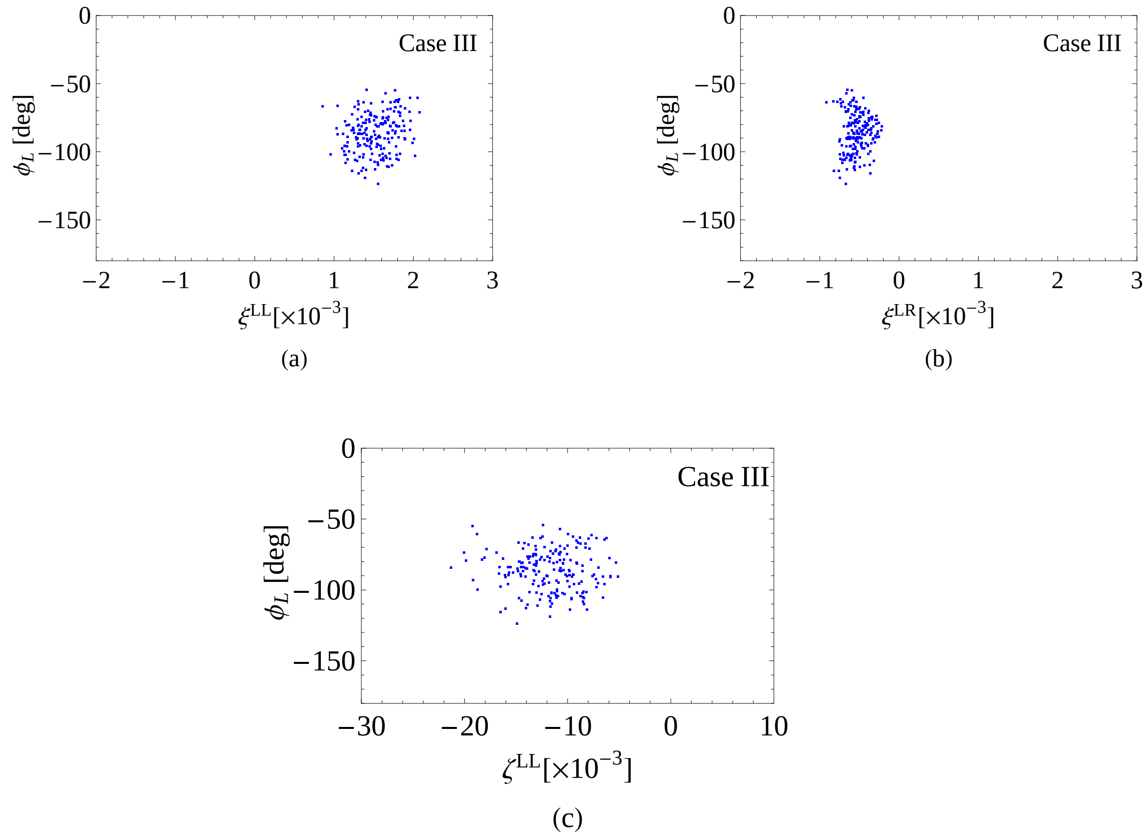

Figure 5: The allowed regions for the parameters $\xi^{LL,LR}$, ζ^{LL}, and ϕ_L in Case III.

with the experimental data within 2σ.

Case IV: Without any simplification of $B_{uu}^{L,R}$ and $B_{dd}^{L,R}$, *i.e.*, arbitrary values of $\zeta^{LL,LR}$ and $\xi^{LL,LR}$ are allowed.

More generally, we give up any assumptions of the couplings $B_{uu}^{L,R}$ and $B_{dd}^{L,R}$. Then, there are five arbitrary NP parameters. As in Case III, we take $\mathcal{B}(B \to \pi K)$, $A_{CP}(B \to \pi K)$ and $A_{CP}(B \to \rho K)$ as constraints and present the predictions for the other observables.

The allowed regions for $\xi^{LL,LR}$, $\zeta^{LL,LR}$ and ϕ_L are shown in Fig. 6, while the numerical results are listed in the last column of Table 5. We find that, similar to Case III, the values of $\xi^{LL,LR}$ are definitely nonzero. The values of ζ^{LL} is a little larger than the one in Case III, due to the interference effect caused by the parameter ζ^{LR}. Our predictions for $\mathcal{B}(B \to \pi K^*, \rho K)$, $A_{CP}(B \to \pi K^*)$ and $A_{CP}^{mix}(B^0 \to \pi^0 K_S, \rho^0 K_S)$, listed in Tables 2, 3 and 4, respectively, are

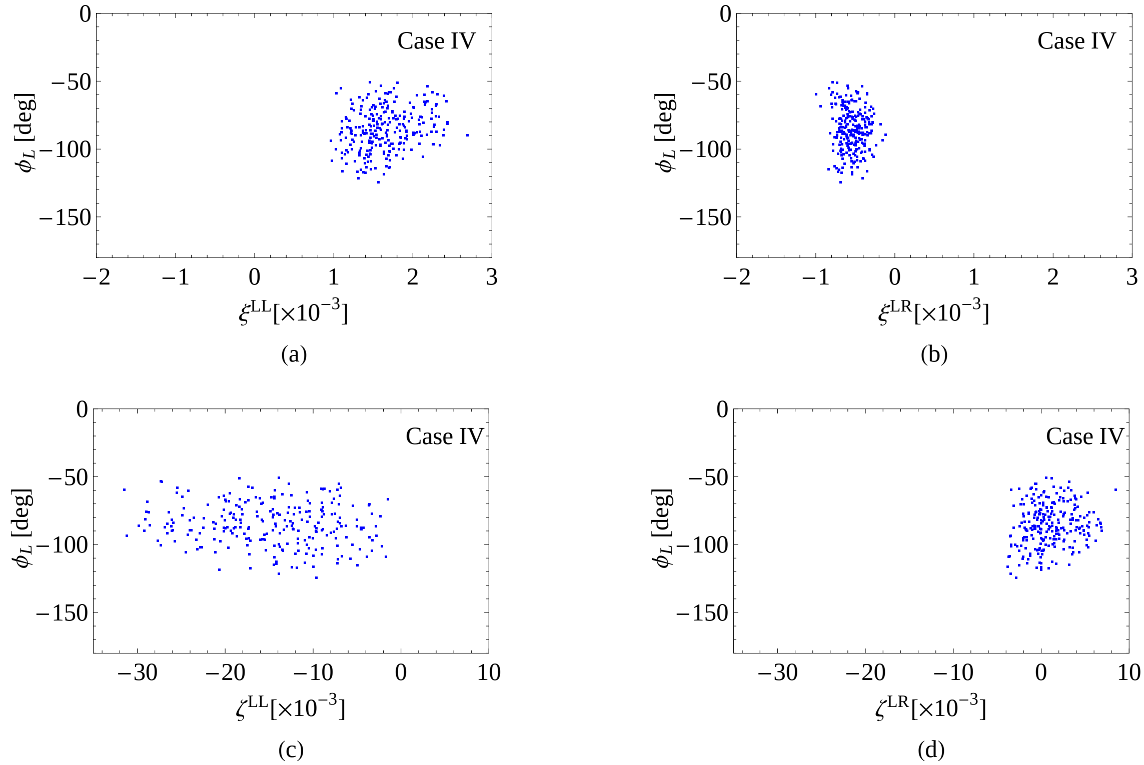

Figure 6: The allowed regions for the parameters $\xi^{LL,LR}$, $\zeta^{LL,LR}$ and ϕ_L in Case IV.

consistent with the experimental data within 2σ.

4 Conclusions

Motivated by the recent observed large difference ΔA between $\mathcal{A}_{CP}(B^{\mp} \to \pi^0 K^{\mp})$ and $A_{CP}(B^0 \to K^{\pm}\pi^{\mp})$, we have investigated the effect of family non-universal Z' model and pursued possible solutions to the observed "πK puzzle". Moreover, we have also taken into account the constraints from the $B \to \pi K^*$, ρK decays, which also involve the same quark level $b \to s\bar{q}q$ $(q = u, d)$ transitions. Our main conclusions are summarized as:

- The Z' contributions to the coefficients of operators O_7 and O_9 (ξ^{LL} and ξ^{LR}) with $\phi_L \sim -86^\circ$ are crucial to bridge the discrepancy of $A_{CP}(B^- \to \pi^0 K^-)$ between theoretical prediction and experimental data. However, they are definitely unequal and opposite in sign.

- The Z' contributions to the coefficients of QCD penguins operator O_3 related to ζ_{LL} are required to moderate the contradiction of $A_{CP}(B^- \to \rho^0 K^-)$ and $A_{CP}(B^0 \to \pi^+ K^-)$ to thier experimental values induced by ξ^{LL} and ξ^{LR}, respectively, even though they are helpless to resolve the observed "πK puzzle". On the other hand, the Z' contributions to $C_5(\zeta^{LR})$ are inessential.

- For all of the four cases, a new weak phase associated with the chiral Z' couplings, with a value about $-86°$, is always required for the "πK puzzle".

Combing the up-to-date experimental measurements of $B \to \pi K$, piK^* and ρK decays, the family non-universal Z' model is found to be helpful to resolve the observed"πK puzzle". It is also reminded that more refined measurements of the mix-induced CP asymmetries in the $B^0 \to \pi^0 K_S$ and $\rho^0 K_S$ decays are required to confirm or refute the NP signals. In the following years, the precision of measurements for these observables is expected to be much improved, which will then shrink and reveal the Z' parameter spaces.

Note added: When the paper is finished, we are aware of the interesting paper by Barger et al.[56]. Although our topics are very similar, we have taken into account of not only the CP asymmetries but also the branching ratios of the correlated decay modes to constrain Z' couplings. Moreover, our approaches for the hadronic dynamics are different.

Acknowledgments

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Appendix A: decay amplitudes in the SM with QCDF

The decay amplitudes for $B \to \pi K$ decays are recapitulated from Ref. [3]

$$\mathcal{A}^{\rm SM}_{B^- \to \pi^- \bar{K}} = \sum_{p=u,c} V_{pb} V^*_{ps} A_{\pi \bar{K}} \Big[\delta_{pu} \beta_2 + \alpha_4^p - \frac{1}{2} \alpha^p_{4,\rm EW} + \beta_3^p + \beta^p_{3,\rm EW} \Big], \tag{24}$$

$$\sqrt{2}\,\mathcal{A}^{\rm SM}_{B^-\to\pi^0K^-} = \sum_{p=u,c} V_{pb}V^*_{ps}\Big\{A_{\pi^0K^-}\Big[\delta_{pu}\,(\alpha_1+\beta_2)+\alpha_4^p+\alpha^p_{4,\rm EW}+\beta_3^p+\beta^p_{3,\rm EW}\Big] + A_{K^-\pi^0}\Big[\delta_{pu}\,\alpha_2+\frac{3}{2}\alpha^p_{3,\rm EW}\Big]\Big\}, \qquad (25)$$

$$\mathcal{A}^{\rm SM}_{\bar{B}^0\to\pi^+K^-} = \sum_{p=u,c} V_{pb}V^*_{ps}A_{\pi^+K^-}\Big[\delta_{pu}\,\alpha_1+\alpha_4^p+\alpha^p_{4,\rm EW}+\beta_3^p-\frac{1}{2}\beta^p_{3,\rm EW}\Big], \qquad (26)$$

$$\sqrt{2}\,\mathcal{A}^{\rm SM}_{\bar{B}^0\to\pi^0\bar{K}^0} = \sum_{p=u,c} V_{pb}V^*_{ps}\Big\{A_{\pi^0\bar{K}^0}\Big[-\alpha_4^p+\frac{1}{2}\alpha^p_{4,\rm EW}-\beta_3^p+\frac{1}{2}\beta^p_{3,\rm EW}\Big] + A_{\bar{K}^0\pi^0}\Big[\delta_{pu}\,\alpha_2+\frac{3}{2}\alpha^p_{3,\rm EW}\Big]\Big\}, \qquad (27)$$

where the explicit expressions for the coefficients $\alpha_i^p \equiv \alpha_i^p(M_1M_2)$ and $\beta_i^p \equiv \beta_i^p(M_1M_2)$ can also be found in Ref. [3]. Note that expressions of the hard-spectator terms H_i appearing in α_i^p and the weak annihilation ones appearing in β_i^p should be replaced by our recalculated ones listed in Appendix B. The decay amplitudes of $B\to\pi K^*$ and $B\to\rho K$ decays could be obtained from the above results by replacing $(\pi K)\to(\pi K^*)$ and $(\pi K)\to(\rho K)$, respectively.

Appendix B: The hard-spectator and annihilation corrections with the infrared finite gluon propagator

With the infrared finite gluon propagator to cure the end-point divergences, the hard-spectator corrections in $B\to PP$ and PV decays can be expressed as [8]

$$H_i(M_1M_2) = \frac{B_{M_1M_2}}{A_{M_1M_2}}\int_0^1 dxdyd\xi\frac{\alpha_s(q^2)}{\xi}\Phi_{B1}(\xi)\Phi_{M_2}(x)\Big[\frac{\Phi_{M_1}(y)}{\bar{x}(\bar{y}+\omega^2(q^2)/\xi)}+r_\chi^{M_1}\frac{\phi_{m_1}(y)}{x(\bar{y}+\omega^2(q^2)/\xi)}\Big], \qquad (28)$$

for the insertion of operators $Q_{i=1-4,9,10}$,

$$H_i(M_1M_2) = -\frac{B_{M_1M_2}}{A_{M_1M_2}}\int_0^1 dxdyd\xi\frac{\alpha_s(q^2)}{\xi}\Phi_{B1}(\xi)\Phi_{M_2}(x)\Big[\frac{\Phi_{M_1}(y)}{x(\bar{y}+\omega^2(q^2)/\xi)}+r_\chi^{M_1}\frac{\phi_{m_1}(y)}{\bar{x}(\bar{y}+\omega^2(q^2)/\xi)}\Big], \qquad (29)$$

for $Q_{i=5,7}$, and $H_i(M_1M_2)=0$ for $Q_{i=6,8}$. When both M_1 and M_2 are pseudoscalars, the final building blocks for annihilation contributions can be expressed as [8]

$$A_1^i = \pi\int_0^1 dxdy\alpha_s(q^2)\Big\{\Big[\frac{\bar{x}}{(\bar{x}y-\omega^2(q^2)+i\epsilon)(1-x\bar{y})}+\frac{1}{(\bar{x}y-\omega^2(q^2)+i\epsilon)\bar{x}}\Big]\Phi_{M_1}(y)\Phi_{M_2}(x) + \frac{2}{\bar{x}y-\omega^2(q^2)+i\epsilon}r_\chi^{M_1}r_\chi^{M_2}\phi_{m_1}(y)\phi_{m_2}(x)\Big\}, \qquad (30)$$

$$A_2^i = \pi\int_0^1 dxdy\alpha_s(q^2)\Big\{\Big[\frac{y}{(\bar{x}y-\omega^2(q^2)+i\epsilon)(1-x\bar{y})}+\frac{1}{(\bar{x}y-\omega^2(q^2)+i\epsilon)y}\Big]\Phi_{M_1}(y)\Phi_{M_2}(x)$$
$$+\frac{2}{\bar{x}y-\omega^2(q^2)+i\epsilon}r_\chi^{M_1}r_\chi^{M_2}\phi_{m_1}(y)\phi_{m_2}(x)\Big\}, \tag{31}$$

$$A_3^i = \pi\int_0^1 dxdy\alpha_s(q^2)\Big\{\frac{2\bar{y}}{(\bar{x}y-\omega^2(q^2)+i\epsilon)(1-x\bar{y})}r_\chi^{M_1}\phi_{m_1}(y)\Phi_{M_2}(x)$$
$$-\frac{2x}{(\bar{x}y-\omega^2(q^2)+i\epsilon)(1-x\bar{y})}r_\chi^{M_2}(x)\phi_{m_2}(x)\Phi_{M_1}(y)\Big\}, \tag{32}$$

$$A_1^f = A_2^f = 0, \tag{33}$$

$$A_3^f = \pi\int_0^1 dxdy\alpha_s(q^2)\Big\{\frac{2(1+\bar{x})}{(\bar{x}y-\omega^2(q^2)+i\epsilon)\bar{x}}r_\chi^{M_1}\phi_{m_1}(y)\Phi_{M_2}(x)$$
$$+\frac{2(1+y)}{(\bar{x}y-\omega^2(q^2)+i\epsilon)y}r_\chi^{M_2}(x)\phi_{m_2}(x)\Phi_{M_1}(y)\Big\}. \tag{34}$$

When M_1 is a vector meson and M_2 a pseudoscalar, the sign of the second term in A_1^i, the first term in A_2^i, and the second terms in A_3^i and A_3^f need to be changed. When M_2 is a vector meson and M_1 a pseudoscalar, one only has to change the overall sign of A_2^i.

Appendix C: Theoretical input parameters

C1. CKM matrix elements

For the CKM matrix elements, we adopt the Wolfenstein parameterization [46] and choose the four parameters A, λ, ρ and η as [47]

$$A = 0.798^{+0.023}_{-0.017}, \quad \lambda = 0.22521^{+0.00083}_{-0.00082}, \quad \overline{\rho} = 0.141^{+0.035}_{-0.021}, \quad \overline{\eta} = 0.340\pm 0.016, \tag{35}$$

with $\overline{\rho} = \rho\,(1-\frac{\lambda^2}{2})$ and $\bar{\eta} = \eta\,(1-\frac{\lambda^2}{2})$.

C2. Quark masses and lifetimes

As for the quark masses, there are two different classes appearing in our calculation. One type is the current quark mass which appears in the factor r_χ^M through the equation of motion for

quarks. This type of quark masses is scale dependent and denoted by $\overline{m}_q$. Here we take

$$\overline{m}_s(\mu)/\overline{m}_q(\mu) = 27.4\pm 0.4 \;\; [48], \quad \overline{m}_s(2\,\mathrm{GeV}) = 87\pm 6\,\mathrm{MeV} \;\; [48], \quad \overline{m}_b(\overline{m}_b) = 4.20^{+0.17}_{-0.07}\,\mathrm{GeV}\;[49]\,, \tag{36}$$

where $\overline{m}_q(\mu) = (\overline{m}_u + \overline{m}_d)(\mu)/2$, and the difference between u and d quark is not distinguished.

The other one is the pole quark mass appearing in the evaluation of penguin loop corrections, and denoted by m_q. In this paper, we take [49]

$$m_u = m_d = m_s = 0, \quad m_c = 1.61^{+0.08}_{-0.12}\,\mathrm{GeV}, \quad m_b = 4.79^{+0.19}_{-0.08}\,\mathrm{GeV}. \tag{37}$$

As for the B-meson lifetimes, we take [49] $\tau_{B_u} = 1.638\,\mathrm{ps}$ and $\tau_{B_d} = 1.530\,\mathrm{ps}$, respectively.

C3. The decay constants and form factors

In this paper, we take the heavy-to-light transition form factors [51]

$$\begin{aligned}
&F_0^{B\to\pi}(0) = 0.258\pm 0.031, \quad F_0^{B\to K}(0) = 0.331\pm 0.041, \quad V^{B\to K^*}(0) = 0.411\pm 0.033,\\
&A_0^{B\to K^*}(0) = 0.374\pm 0.034, \quad A_1^{B\to K^*}(0) = 0.292\pm 0.028, \quad V^{B\to\rho}(0) = 0.323\pm 0.030,\\
&A_0^{B\to\rho}(0) = 0.303\pm 0.029, \quad A_1^{B\to\rho}(0) = 0.242\pm 0.023.
\end{aligned} \tag{38}$$

and the decay constants

$$\begin{aligned}
&f_B = (216\pm 22)\ \mathrm{MeV}\ [50]\,, \quad f_\pi = (130.4\pm 0.2)\ \mathrm{MeV}\ [49]\,, \quad f_K = (155.5\pm 0.8)\ \mathrm{MeV}\ [49]\,,\\
&f_{K^*} = (217\pm 5)\ \mathrm{MeV}\ [51], \quad f_\rho = (209\pm 2)\ \mathrm{MeV}\ [51].
\end{aligned} \tag{39}$$

C4. The LCDAs of mesons and light-cone projector operators.

The light-cone projector operators of light mesons in momentum space read [52, 3]

$$M^P_{\alpha\beta} = \frac{if_P}{4}\left[\not{p}\,\gamma_5\,\Phi_P(x) - \mu_P\gamma_5\,\frac{\not{k}_2\,\not{k}_1}{k_2\cdot k_1}\,\Phi_p(x)\right]_{\alpha\beta}\,, \tag{40}$$

$$\left(M^V_{\parallel}\right)_{\alpha\beta} = -\frac{if_V}{4}\left[\not{p}\,\Phi_V(x) - \frac{m_V f_V^{\perp}}{f_V}\,\frac{\not{k}_2\,\not{k}_1}{k_2\cdot k_1}\,\Phi_v(x)\right]_{\alpha\beta}\,, \tag{41}$$

where $f_{P,V}$ are the decay constants, and $\mu_P = m_b r_\chi^P/2$, with the chirally-enhanced factor r_χ^P defined as

$$r_\chi^\pi(\mu) \;=\; \frac{2m_\pi^2}{m_b(\mu)2m_q(\mu)}\,, \quad r_\chi^K(\mu) = \frac{2m_K^2}{m_b(\mu)(m_q+m_s)(\mu)}\,, \tag{42}$$

where the quark masses are all running masses defined in the $\overline{\mathrm{MS}}$ scheme. For the LCDAs of mesons, we use their asymptotic forms [53, 54]

$$\Phi_{P,V}(x) = 6\,x(1-x)\,, \quad \phi_p(x) = 1\,, \quad \phi_v(x) = 3(2x-1)\,. \tag{43}$$

As for the B-meson wave function, we take the form [55]

$$\Phi_B(\xi) = N_B \xi(1-\xi)\mathrm{exp}\Big[-\Big(\frac{M_B}{M_B - m_b}\Big)^2(\xi-\xi_B)^2\Big], \tag{44}$$

where $\xi_B \equiv 1 - m_b/M_B$, and N_B is the normalization constant to insure that $\int_0^1 d\xi \Phi_B(\xi) = 1$.

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